

# The compact star in the SU(3) sigma model

Ryszard Manka, Ilona Bednarek

University of Silesia, Katowice, Poland.

## Abstract

The linear chiral  $SU(3)_L \times SU(3)_R$  model is applied to describe properties of the compact star matter inside the quark, protoneutron and neutron star.

## 1 Introduction

In this presentation we shall investigate the astrophysical properties of the compact stars (protoneutron, neutron and quark stars [1]) within the Relativistic Mean Field (RMF) [2] theory originated from the linear SU(3) sigma model [3, 4, 5].

The main aim of this work is to show how an effective mean field approximation (RMF) emerges from the linear SU(3) chiral model and to show comparison to the current RMF approach (Furnstahl - Serot - Tang (FST) model [6]) as well as to present astrophysical applications in the form of the neutron star model. The effective model which includes scalar, vector and scalar-vector interaction terms is applied to describe properties of the quark, protoneutron and neutron star matter.

## 2 The SU(3) sigma model

The chiral SU(3) model was proposed by Papazoglou *et al.* [3, 4]. In the original form it describes interaction of the baryons and mesons SU(3) multiplets. Recently, a chiral SU(3) quark model has been proposed by Wang

*et al* [7]. The basic fields that compose the theory represent the realization of the group  $SU(3)_L \times SU(3)_R$ . The meson content of the model is scalar, pseudo-scalar and vector. Naive quark models interpret them as excited  $\bar{q}q$  states. Scalar and pseudo-scalar mesons can be grouped into

$$\Phi = \Sigma + i\Pi = \frac{1}{\sqrt{2}}T_a\phi_a = \frac{1}{\sqrt{2}}T_a(\sigma_a + i\pi_a) \quad (1)$$

where  $\sigma_a$  and  $\pi_a$  are members of the scalar and pseudo-scalar octet respectively:

$$\Sigma = \frac{1}{\sqrt{2}}\sigma^a\lambda^a = \begin{pmatrix} \frac{1}{\sqrt{2}}(f_0 + a_0^0), & a_0^+, & K^+ \\ a_0^-, & \frac{1}{\sqrt{2}}(f_0 - a_0^0), & K^0 \\ K^-, & \bar{K}^0, & f_0' \end{pmatrix} \quad (2)$$

$$\Pi = \frac{1}{\sqrt{2}}\pi^a\lambda^a = \begin{pmatrix} \frac{1}{\sqrt{2}}(\pi^0 + \eta^8) & \pi^+ & \kappa^+ \\ \pi^- & \frac{1}{\sqrt{2}}(-\pi^0 + \eta^8) & \kappa^0 \\ \kappa^- & \kappa^0 & \zeta \end{pmatrix}. \quad (3)$$

The vector meson octet is given by

$$V = \frac{1}{\sqrt{2}}v^a\lambda^a = \begin{pmatrix} \frac{1}{\sqrt{2}}(\omega + \rho_0^0), & \rho_0^+, & K^{*+} \\ \rho_0^-, & \frac{1}{\sqrt{2}}(\omega - \rho_0^0), & K^{*0} \\ K^{*-}, & \bar{K}^{*0}, & \phi \end{pmatrix}. \quad (4)$$

The most general form of the Lagrangian function can be written as a sum of the following parts

$$\mathcal{L} = \mathcal{L}_M + \mathcal{L}_F + \mathcal{L}_{sb} \quad (5)$$

where

$$\mathcal{L}_M = \frac{1}{2}Tr(\partial_\mu\Phi\partial^\mu\Phi) - \frac{1}{2}\mu^2Tr(\Phi^2) - \frac{\lambda}{4}Tr((\Phi^+\Phi)^2) - \frac{\kappa}{4}Tr(\Phi^+\Phi)^2 \quad (6)$$

is the Lagrangian function which describes scalar and pseudo-scalar mesons  $\Phi$  ( $\Phi = \Sigma + i\Pi$ ). The symmetry breaking term  $\mathcal{L}_{sb}$  has the form of:

$$\mathcal{L}_{sb} = \frac{1}{2}c(Det(\Phi) + Det(\Phi)^*) + Tr(H^+\Phi + \Phi^+H). \quad (7)$$

In the mean field approximation the chiral symmetry is broken and the meson fields gain non-vanishing vacuum expectation values  $(\sigma, \chi)$

$$\langle \Phi \rangle = \langle \Sigma \rangle = \begin{pmatrix} \frac{1}{\sqrt{2}}\sigma, & 0, & 0 \\ a_0^-, & \frac{1}{\sqrt{2}}\sigma, & 0 \\ 0, & 0, & \chi \end{pmatrix}.$$

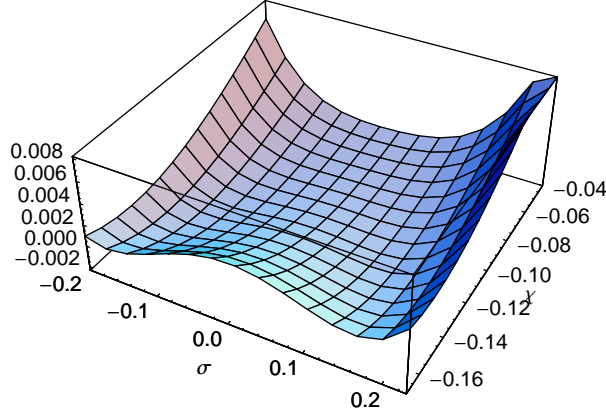


Figure 1: The effective potential  $U_{eff}(\sigma, \chi)$  in the SU(3) sigma model.

The effective potential  $U_{eff}(\sigma, \chi) = - \langle L \rangle$  (Fig. 1) determines the scale of the chiral symmetry breaking. Shifting the meson field  $\Phi = \bar{\Phi} + \langle \Phi \rangle$  ( $f_0 = \bar{f}_0 + \sigma$ ,  $f'_0 = \bar{f}'_0 + \chi$ ) and diagonalizing the square mass matrices  $m_{a,b}^2 = \frac{\partial^2 U_{eff}}{\partial \bar{\sigma}_a \partial \bar{\sigma}_b}$  produce the physical meson fields

$$\begin{Bmatrix} \varphi \\ \varphi_* \end{Bmatrix} = \begin{pmatrix} \cos\vartheta & \sin\vartheta \\ -\sin\vartheta & \cos\vartheta \end{pmatrix} \begin{Bmatrix} f_0 \\ f'_0 \end{Bmatrix}. \quad (8)$$

Fitting the model to the observed masses of mesons allows to determine its parameters (similar to the case [8] of the explicit chiral symmetry breaking with  $U(1)_A$  anomaly). It gives  $\mu = 552.274 \text{ MeV}$ ,  $\lambda = 45.11$ ,  $\kappa = -8.917$  and  $c = 3412.25$ . This fitting gives meson masses (e.g.  $\pi$ ,  $K$ ,  $\varphi$ ,  $\varphi_*$ ,  $\delta = a_0$ , etc.) including sigma meson mass  $m_\sigma = 502.27 \text{ MeV}$ . The only unclear thing is the sigma meson mass. The light scalar meson  $\sigma$  (denoted here as  $\varphi$ ) is an elusive subject of classification.

The fermion contents of the model consists of either quarks or baryons, respectively (QMF or RMF model).

The chiral SU(3) quark mean field model has been applied to describe the quark matter or nucleon matter. In the chiral limit, the quark field  $q = \{u, d, s\}$  with three flavors, can be decomposed into left and right-handed

parts  $q = q_L + q_R$ . The quarks are described by the Lagrange function

$$\mathcal{L}_F = i\bar{q}\gamma^\mu D_\mu q - \bar{q}m_0 q + g_s \bar{q}\Phi q - \chi_c(r)\bar{q}q$$

where  $\chi_c(r)$  is quarks confining potential. Solving the Dirac equation for quark in confining potential the baryon masses can be calculated

$$M_{eff,N}(\varphi, \varphi_*) = M_N - g_\sigma(\varphi) \varphi - g_{\sigma*}\varphi_* = M_N - g_\sigma \varphi + \frac{1}{2}g_\sigma C'(0) \varphi^2 + \dots$$

with  $g_\sigma(\varphi) = g_\sigma - \frac{1}{2}g_\sigma C'(0) \varphi = g_\sigma - \frac{1}{2}a \varphi$ . The last nonlinear term indicates the inner nucleon structure.

### 3 The effective RMF approach

The nuclear relativistic mean field approach describes the nuclear interactions due to the mesons exchange between baryons ( $p, n, \Lambda, \Sigma, \Xi$ ). Baryons are grouped into the isospin and hipercharge representations  $(\frac{1}{2}, 1), (\frac{1}{2}, -1), (1, 0)$

$$\Lambda, N = \begin{pmatrix} p \\ n \end{pmatrix}, \Xi = \begin{pmatrix} \Xi^0 \\ \Xi^- \end{pmatrix}, \Sigma = \begin{pmatrix} \Sigma^+ \\ \Sigma^0 \\ \Sigma^- \end{pmatrix}. \quad (9)$$

The RMF Lagrange functions

$$\mathcal{L}_{RMF} = \mathcal{L}_B + \mathcal{L}_M \quad (10)$$

describes baryons ( $B = \{b, n, \Lambda, \Sigma, \Xi\}$ )

$$\mathcal{L}_B = \sum_B i\bar{\psi}_B \gamma^\mu D_\mu \psi_B - \sum_B M_B(\varphi, \varphi_*) \bar{\psi}_B \psi_B \quad (11)$$

and mesons

$$\mathcal{L}_M = \mathcal{L}_{Ms} + \mathcal{L}_{Mv}. \quad (12)$$

Mesons can be divided into scalar mesons ( $\varphi, \varphi_*, \delta$ ) described by  $\mathcal{L}_{Ms}$  and vector mesons ( $\omega, \rho, \phi$ ) described by  $\mathcal{L}_{Mv}$ .

$$\begin{aligned} \mathcal{L}_{Ms} = & \frac{1}{2}\partial_\mu \varphi \partial^\mu \varphi + \frac{1}{2}\partial_\mu \delta^a \partial^\mu \delta^a + + \\ & \frac{1}{2}\partial_\mu \varphi_* \partial^\mu \varphi_* - U_{S,eff}(\varphi, \varphi_*, \delta), \end{aligned} \quad (13)$$

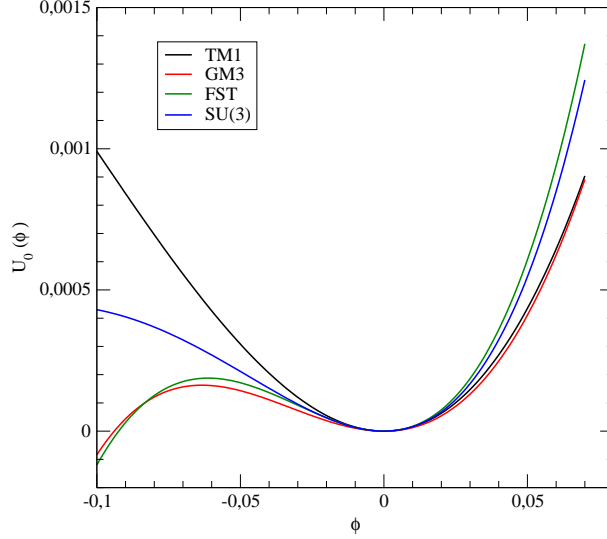


Figure 2: The potential  $U_0$  for the scalar meson  $\varphi$  of the effective RMF theory.

$$\begin{aligned} \mathcal{L}_{Mw} = & -\frac{1}{4}\Omega_{\mu\nu}\Omega^{\mu\nu} + \frac{1}{2}m_\omega^2\omega_\mu\omega^\mu - \frac{1}{4}R_{\mu\nu}^a R^{a\mu\nu} + \frac{1}{2}m_\rho^2\rho_\mu^a\rho^{a\mu} \\ & -\frac{1}{4}\Phi_{\mu\nu}\Phi^{\mu\nu} + \frac{1}{2}m_\phi^2\phi_\mu\phi^\mu + U_{V,eff}(\omega, \rho) \end{aligned} \quad (14)$$

where

$$\Omega_{\mu\nu} = \partial_\mu\omega_\nu - \partial_\nu\omega_\mu \quad , \quad \Phi_{\mu\nu} = \partial_\mu\phi_\nu - \partial_\nu\phi_\mu \quad (15)$$

$$R_{\mu\nu} = R_{\mu\nu}^a T^a = \partial_\mu R_\nu - \partial_\nu R_\mu - ig_\rho[R_\mu, R_\nu]. \quad (16)$$

The scalar meson interaction is enormously nonlinear. It comes from the shifting in the potential  $U_{eff}(\sigma, \chi)$  according to the prescription (eq. 8). In the simplest approximation this procedure generates the polynomial scalar interaction (Fig. 2) of the RMF approach

$$U_0(\varphi) = U_{eff}(\sigma_0 + \varphi, \chi_0) = U_{eff}(\varphi, 0) = \frac{1}{2}m_\sigma^2\varphi^2 + \frac{1}{3}g_2\varphi^3 + \frac{1}{4}g_3\varphi^4 \quad (17)$$

The form of the potential function was first introduced by Boguta and Bodmer [9] in order to get the correct value of the compressibility  $K$  of nuclear matter at saturation density (see Table 1). The simplest Walecka model (L2) (*linear* Walecka model  $g_2 = g_3 = 0$ ) brings a very large, unrealistic value

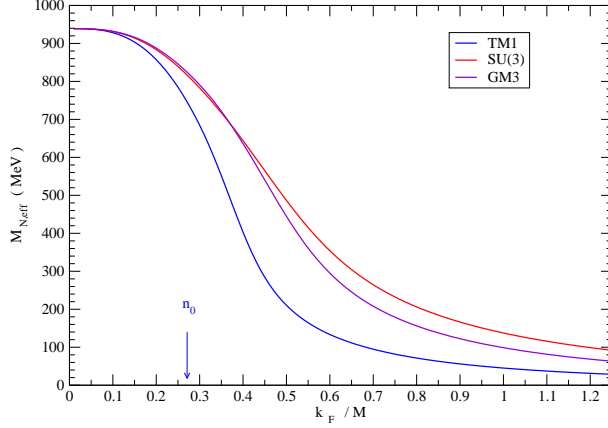


Figure 3: The effective nucleon masses for different parameter sets as a function of the nucleon Fermi momentum  $k_F$ .

Parameter	GM3[10]	TM1[11]	FST[6]	SU(3)
$E_0$ (MeV)	-16.35	-16.26	-16.38	-16.31
$\delta_0$	0.793	0.659	0.661	0.761
$n_0$ ( $fm^{-3}$ )	0.153	0.145	0.155	0.145
$K$ (MeV)	241.12	281.53	219.5	194.9
$J$ (MeV)	32.44	36.82	38.17	34.70

Table 1: Properties of the nuclear matter at saturation for the symmetric nuclear matter.

of the parameter  $K$  [2]. Fig. 3 depicts the effective nucleon masses obtained for different parameter sets functions of baryon number density  $n_B$ . The parameters describing the nucleon-nucleon interactions in the RMF approach are chosen in order to reproduce the properties of the symmetric nuclear matter at saturation such as the binding energy, symmetry energy and incompressibility. In the chiral  $SU(3)$  model they are generally calculable from the starting ones ( $\mu^2$ ,  $\lambda$ ,  $\kappa$ ) which are fitted to the mesons spectroscopy. The appropriate parameter set is constrained not only by the value of physical scalar meson masses but also by the properties of nuclear matter at saturation. For symmetric nuclear matter the nucleon density equals  $n_0 = 2.5 \cdot 10^{14} \text{ g cm}^{-3} = 0.15 \text{ nucleons/fm}^3 = 140 \text{ MeV fm}^{-3}$ . The obtained results are collected in Table 1.

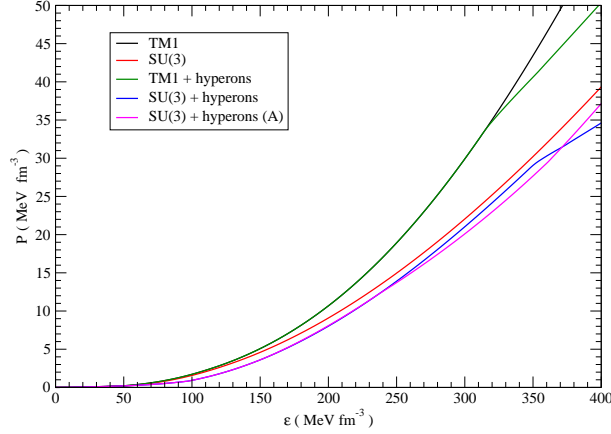


Figure 4: The nuclear matter equation of state for different parameter sets.

## 4 The compact star in the SU(3) sigma model

The compact star is a result of the equilibrium between gravitational collapse and the pressure generated by the nuclear or quark matter and leptons [12, 13]. With respect to the electroweak interaction the matter is in  $\beta$  equilibrium. In the neutron star weak interactions are responsible for  $\beta$  decay

$$n + \nu_e \leftrightarrow p + e, \quad (18)$$

$$\mu + \nu_e \leftrightarrow e + \nu_\mu. \quad (19)$$

These reactions produce appropriate relation among the chemical potentials of neutrinos in a protoneutron star where they are trapped

$$\mu_p = \mu_n + \mu_{\nu_e} - \mu_e \quad (20)$$

$$\mu_{\nu_e} = \mu_e + \mu_p - \mu_n \quad (21)$$

$$\mu_{\nu_\mu} = \mu_{\nu_e} - \mu_e + \mu_\mu \quad . \quad (22)$$

### 4.1 The protoneutron star.

Protoneutron stars are hot and lepton rich objects formed as a result of type II supernovae explosion. The collapse of an iron core of a massive star leads to the formation of a core residue which is considered as an intermediate stage before the formation of a cold, compact neutron star. This intermediate stage

which is called a protoneutron star can be described [19, 20, 21, 22] as a hot, neutrino opaque core, surrounded by a colder neutrino transparent outer envelope. The evolution of a nascent neutron star can be described by a series of separate phases starting from the moment when the star becomes gravitationally decoupled from the expanding ejecta. In this paper two evolutionary phases which can be characterized by the following assumptions:

- the low entropy core  $s/n_B = 1-2$  (in units of the Boltzmann's constant) with trapped neutrinos  $Y_L = 0.4$
- the cold, deleptonized core ( $Y_L = 0, s/n_B = 0$ ).

have been considered. These two distinct stages are separated by the period of deleptonization. During this epoch the neutrino fraction decreases from the nonzero initial value ( $Y_\nu \neq 0$ ) which is established by the requirement of the fixed total lepton number at  $Y_l = 0.4$ , to the final one characterized by  $Y_\nu = 0$ . Evolution of a protoneutron star which proceeds by neutrino emission causes that the star changes itself from a hot, bloated object to a cold, compact neutron star.

The interior of this very early stage of a protoneutron star is an environment in which matter with the value of entropy of the order of 2 with trapped neutrinos produces a pressure to oppose gravitational collapse. The lepton composition of matter is specified by the fixed lepton number  $Y_L = 0.4$ . Conditions that are indispensable for the unique determination of the equilibrium composition of a protoneutron star matter arise from the requirement of  $\beta$  equilibrium, charge neutrality and baryon and lepton number conservation. The later one is strictly connected with the assumption that the net neutrino fraction  $Y_\nu \neq 0$  and therefore the neutrino chemical potential  $\mu_\nu \neq 0$ . When the electron chemical potential  $\mu_e$  reaches the value equal to the muon mass, muons start to appear.

In a protoneutron star, when the neutrino is trapped, its chemical potential strongly depends on nuclear asymmetry

$$\mu_{\nu_e} = \mu_e - \mu_A, \quad (23)$$

$$\mu_A = \mu_n - \mu_p = \varepsilon_n - \varepsilon_p - g_\rho r_0 \quad (24)$$

$$\epsilon_{p,n} = \sqrt{k_{F,B}^2 + M_{B,eff}^2}|_{B=n,p}. \quad (25)$$



where  $r_0 = \langle \rho_0^3 \rangle$  is the expected value for the  $\rho$  meson in medium. If  $\nu_\mu$  are able to escape from the protoneutron star and  $\mu_e$  not, the muon chemical potential

$$\mu_\mu = \mu_A$$

depends on the nuclear asymmetry. In conclusion, the protoneutron star with the biggest neutrino number is in the nuclear symmetric phase, when  $\mu_A = 0$ , so  $\mu_{\nu_e} = \mu_e$  and muons are absent  $\mu_\mu = 0$ . The neutrinos as fermions increase the star pressure and make the star with a large radius ( $\sim 20-50 \text{ km}$ , Fig. 5). It is interesting that the outer layers when density drops below  $\sim 10^{14} \text{ g/cm}^3$  and nuclear interaction vanishes the weak interactions with  $\beta$  decay still takes place as long as neutrinos are trapped. The outer layers of the protoneutron star has mainly the electroweak nature. The star core is similar to the symmetric hot nuclear matter. The entropy per baryon  $s/n_B = 1$  gives temperature ( $\sim 40-80 \text{ MeV}$ ) inside the star.

## 4.2 The neutron star.

When neutrinos escape at last ( $\mu_{\nu_e} = 0$ ), the new equilibrium have to take place at least in the neutron star case. Neutrinos are now neglected here since they leak out from the neutron star, whose energy diminishes at the same time. The pressure decrease reduces the star radius till  $\sim 10-14 \text{ km}$  and increases the star density to  $\sim 10^{14} - 10^{15} \text{ g/cm}^3$ , the nuclear interactions became more crucial. The equilibrium conditions with respect to the  $\beta$  decay between baryonic (including hyperons) and leptonic species lead to the following relations among their chemical potentials and constrain the species fraction in the star interior

$$\mu_p = \mu_{\Sigma^+} = \mu_n - \mu_e \quad \mu_\Lambda = \mu_{\Sigma^0} = \mu_{\Xi^0} = \mu_n \quad (26)$$

$$\mu_{\Sigma^-} = \mu_{\Xi^-} = \mu_n + \mu_e \quad \mu_\mu = \mu_e \quad (27)$$

The higher the density the greater number of hyperon species are expected to appear. They can be formed both in leptonic and baryonic processes. In the latter one the strong interaction process such as

$$n + n \rightarrow n + \Lambda \quad (28)$$

proceeds. There are other relevant strong reactions that establish the hadron population in neutron star matter e.g.:

$$\Lambda + n \rightarrow \Sigma^- + p \quad \Lambda + \Lambda \rightarrow \Xi^- + p \quad (29)$$

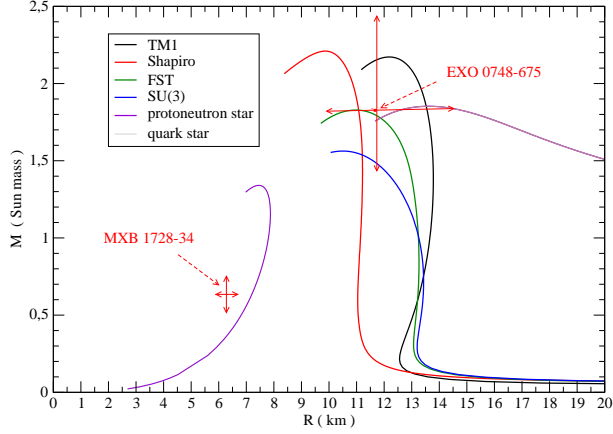


Figure 5: The mass radius relation for various quark and neutron stars.

The final result is the equation of the state (Fig. 4). All these, lead to the neutron star model with the value of maximum mass close to  $1.5 M_{\odot}$  with the reduced value of proton fraction and very compact hyperon core. The FST equation of state gives more massive neutron star  $\sim 1.7 - 1.8 M_{\odot}$ .

### 4.3 The quark star.

Quark strange stars are astrophysical compact objects which are entirely made of deconfined  $u, d, s$  quark matter (*strange matter*) staying in  $\beta$  - equilibrium. The possible existence of strange stars is a direct consequence of the conjecture that strange matter may be the absolute ground state of strongly interacting matter.

Some years ago Guichon proposed an interesting model concerning the change of the nucleon properties in nuclear matter (quark-meson coupling model (QMC)) [14]. The model construction mimics the relativistic mean field theory, where the scalar  $\sigma$  and the vector meson  $\omega$  fields couple not with nucleons but directly with quarks. The quark mass has to change from its bare current mass due to the coupling to the  $\sigma$  meson. More recently, Shen and Toki [15] have proposed a new version of the QMC model, where the interaction takes place between constituent quarks and mesons. They refer the model as the quark mean field model (QMF). In this work we shall also investigate the quark matter within the QMF theory motivated by

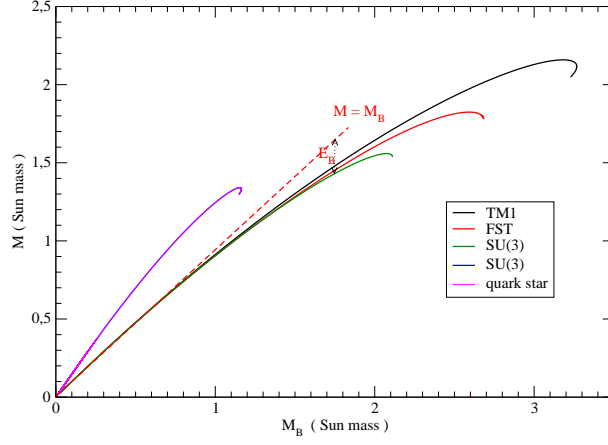


Figure 6: The gravitational binding energy for quark and neutron stars.

parameters coming from the SU(3) chiral model. The quark phase consists of the constituent quarks with three flavors  $u$ ,  $d$  and  $s$ . Quarks are effectively free quasiparticles in vacuum with non-vanishing bag 'constant'. Quarks and electrons are in  $\beta$ -equilibrium which can be described as a relation among their chemical potentials

$$\mu_d = \mu_u + \mu_e = \mu_s$$

where  $\mu_u$ ,  $\mu_d$ ,  $\mu_s$  and  $\mu_e$  stand for quarks and electron chemical potentials respectively. If the electron Fermi energy is high enough (greater than the muon mass) in the neutron star matter muons start to appear as a result of the following reaction

$$\begin{aligned} d &\rightarrow u + e^- + \bar{\nu}_e \\ s &\rightarrow u + \mu^- + \bar{\nu}_\mu \end{aligned}$$

The neutron chemical potential is

$$\mu_n \equiv \mu_u + 2\mu_d.$$

In a pure quark state the star should be charge neutral. This gives us an additional constraint on the chemical potentials

$$\frac{2}{3}n_u - \frac{1}{3}n_d - \frac{1}{3}n_s - n_e = 0. \quad (30)$$

where  $n_f$  ( $f \in u, d, s$ ),  $n_e$  the particle densities of the quarks and the electrons, respectively. The EOS can now be parameterized by only one chemical potential, say  $\mu_u$ . Nowadays the strange quark star is the subject of considerable interest [16, 17, 18].

#### 4.4 The Tolman - Oppenheimer - Volkoff equations.

The spherically symmetric static star in general gravity is solution of the Einstein equations

$$R_{\mu\nu} - \frac{1}{2}g_{\mu\nu}R = \kappa T_{\mu\nu}$$

for the metric

$$g_{\mu\nu} = \begin{pmatrix} -e^{\nu(r)} & 0 & 0 & 0 \\ 0 & e^{\lambda(r)} & 0 & 0 \\ 0 & 0 & r^2 & 0 \\ 0 & 0 & 0 & r^2 \sin^2 \theta \end{pmatrix}$$

where  $\kappa = \frac{8\pi G}{c^4}$  and

$$T_{\mu\nu} = (P + \epsilon)u_\mu u_\nu - P g_{\mu\nu} = \begin{pmatrix} \epsilon = c^2 \rho & 0 & 0 & 0 \\ 0 & P & 0 & 0 \\ 0 & 0 & P & 0 \\ 0 & 0 & 0 & P \end{pmatrix}$$

is the energy-momentum tensor describing an equation of state of the matter. One of the Einstein equations gives

$$e^{-\lambda} = 1 - \frac{2Gm(r)}{r}$$

where  $m(r)$  is the mass accumulated inside a sphere of radius, so

$$\frac{dm(r)}{dr} = 4\pi r^2 \rho. \quad (31)$$

The continuity equation

$$T^{\mu\nu}_{;\nu} = 0 \Rightarrow \frac{d\nu(r)}{dr} = -\frac{2}{P(r) + c^2 \rho(r)} \frac{dP(r)}{dr}.$$

The second Einstein equation gives the Tolman -Oppenheimer-Volkoff equation

$$\frac{dP(r)}{dr} = -\frac{Gm(r)\rho(r)}{r^2} \frac{(1 + \frac{P(r)}{\rho(r)})(1 + \frac{4\pi r^3 P(r)}{m(r)})}{1 - \frac{2Gm(r)}{r}}. \quad (32)$$

The obtained form of the equation of state serves as an input to the Tolman - Oppenheimer - Volkoff equations and determines the structure of spherically symmetric stars. The gravitational binding energy of a relativistic star is defined as a difference between its gravitational and baryon masses

$$E_{b,g} = (M_p - m(R))c^2 \quad (33)$$

where

$$M_p = 4\pi \int_0^R dr r^2 (1 - \frac{2Gm(r)}{c^2 r})^{-\frac{1}{2}} \rho(r). \quad (34)$$

The numerical solution of the above equation is of considerable relevance to the selected EOS. Numerical solutions of these equations allow to construct the mass-radius relations of the neutron star (Fig. 5) and the gravitational final energy (Fig. 6).

## 5 Conclusions

The physics of compact objects like neutron stars offers an intriguing interplay between nuclear processes and astrophysical observables. Neutron stars exhibit conditions far from those encountered on earth. The determination of an equation of state (EoS) for dense matter is essential for calculations of neutron star properties. It is the equation of state that determines the characteristics and properties of neutron stars such as possible range of masses, the mass-radius relation and the cooling rate. The astrophysical properties of the compact stars are strongly related to the properties of quark and nuclear matter. Its mass, radius and binding energy depend on inner physics of an elementary particle and nucleus. The linear sigma model predicts properties of the compact stars which looks quite reasonable. Only the quark star seems to look as an unphysical object (the positive binding energy, Fig. 6). The astronomical observation of a mass - radius relation (for example the SAX J1808.4-3658 [23] compact star) may be used to determine the nuclear equation of state. The best fit to the newly measured radius and mass for EXO 0748-675 neutron star [24] rather prefers (Fig. 5) more traditional FST equation of state [6].

## References

- [1] N. K. Glendenning, *Compact Stars* by N. K Glendenning Springer-Verlag, New York, 1997
- [2] B. D. Serot, J. D. Walecka 1997 *Recent Progress in Quantum Hadrodynamics Int. J. Mod. Phys.* **E6** 515; [nucl-th/9701058](#).
- [3] P. Papazoglou, J. Schramm, J. Schaffer - Bielich, H. Stocker and W. Greiner, Phys. Rev. C. **57** (1998) 2576.
- [4] P. Papazoglou, D. Zschesche, J. Schramm, J. Schaffer - Bielich, H. Stocker and W. Greiner, Phys. Rev. C. **59** (1999) 411.
- [5] N. A. Tornqvist, Eur. Phys.J. **C11** (1999) 359-363; [hep-ph/9905282](#); F.E.Close,N. A. Tornqvist, J Phys. **G28** (2002) R249; [hep-ph/0204205](#);
- [6] B. D. Serot, Lect. Notes Phys 641 (2004) 31; [nucl-th/0308047](#). B. Serot, *Covariant Effective Field Theory for Nuclear Structure and Currents*, [nucl-th/0405058](#);
- [7] P. Wang, Z. Y. Zhang,Y. W. Yu, R. K. Su and H. Q. Song, Nucl. Phys. A. **688** (2001) 791
- [8] D. Roder, J. Ruppert, D. H. Rischke, Phys.Rev. **D68**, (2003) 016003 ;[nucl-th/0301085](#);
- [9] A. R. Bodmer, C. E. Price, Nucl.Phys. **A505**, 123 (1989); A. R. Bodmer, Nucl.Phys. **A526**, 703 (1991).
- [10] N. K. Glendenning, F. Weber, S. A. Moszkowski, Phys. Rev. **C45**, (1992), 844
- [11] Y. Sugahara and H. Toki, Prog. Theo. Phys. **92**, 803 (1994).
- [12] M. Prakash, *The Equation of State and Neutron Star*, lectures delivered at the Winter School held in Puri, India, (1994); J.M. Lattimer, M. Prakash, *Neutron star structure and equation of state* [astro-ph/0002232](#).
- [13] J. Blaizot, D. Gogny and B. Grammaticos, Nucl. Phys. **A265**, 315 (1975).

- [14] P. Guichon 1988 Phys. Lett. **B200**, 235; P. Guichon, K. Saito, E. Rodionov, A.W. Thomas, *The role of nucleon structure in finite nuclei*, [nucl-th/9509034](#)
- [15] H. Shen, H. Toki 1978 Phys. Rev. **D18**, 4187, *Quark mean field model for nuclear matter and finite nuclei*, [nucl-th/9911046](#)
- [16] R. Manka, I. Bednarek, G. Przybyla, New J. Phys. **4** (2002) 14, [nucl-th/0201003](#)
- [17] M. Bubulla, *NJL-model analysis of quark matter at large density*, [hep-ph/0402234](#)
- [18] F. Weber, *Strange Quark Matter and Compact Stars* [astro-ph/0407155](#)
- [19] M. Prakash, I. Bombaci, Manju Prakash, P.J. Ellis, J.M. Lattimer, R. Knorren 1997 *Phys. Rep.* **280** 1
- [20] J.M. Lattimer, M. Prakash 2001 *Astrophys. J.* **550** 426
- [21] J.A. Pons, S. Reddy, P.J. Ellis, M. Prakash, J.M. Lattimer *Phys. Rev.* **C62** 035803
- [22] A. Burrows, J.M. Lattimer 1986 *Astrophys. J.* **307** 178
- [23] X. -D. Li, I. Bombaci, J. Dey, M. Dey, E.P.J. van den Heuvel, Phys. Rev. Lett. **83**, 3776 (1999); J. Dey, S. Ray, X. -D. Li, M. Dey, I. Bombaci, *Glimpses of the strange star*, [astro-ph/0001305](#).
- [24] A.R. Villarreal, T.E. Strohmayer, *Discovery of the Neutron Star Spin Frequency in EXO 0748-675*, [astro-ph/0409384](#)